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Resonant leptogenesis in the minimal $B - L$ extended standard model at TeV

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We investigate the resonant leptogenesis scenario in the minimal $B - L$ extended standard model with the $B - L$ symmetry breaking at the TeV scale. Through detailed analysis of the Boltzmann equations, we show how much the resultant baryon asymmetry via leptogenesis is enhanced or suppressed, depending on the model parameters, in particular, the neutrino Dirac-Yukawa couplings and the TeV scale Majorana masses of heavy degenerate neutrinos. In order to consider a realistic case, we impose a simple ansatz for the model parameters and analyze the neutrino oscillation parameters and the baryon asymmetry via leptogenesis as a function of only a single CP phase. We find that for a fixed CP phase all neutrino oscillation data and the observed baryon asymmetry of the present Universe can be simultaneously reproduced.

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I. INTRODUCTION

The origin of the baryon asymmetry in the present Universe is one of the big mysteries in cosmology. The ratio of the baryon (minus antibaryon) density n_B to the entropy density s has been measured with the precision at 10% level by the Wilkinson Microwave Anisotropy Probe satellite experiment [1],

$$Y_B = \frac{n_B}{s} = 0.87 \times 10^{-10}. \tag{1}$$

It might be the most attractive if the origin of the baryon asymmetry can be explained within the context of the standard model (SM), the electroweak baryogenesis [2]. In order for this scenario to work, a strong first order electroweak phase transition is necessary in the early Universe. However, the Higgs potential satisfying the current lower bound on the SM Higgs boson mass [3] is not likely to show this strong first order phase transition and hence, the SM electroweak baryogenesis is almost ruled out.

An appealing alternative is the leptogenesis scenario [4], which is also intimately related with the smallness of the neutrino masses through the seesaw mechanism [5]. A most widely accepted scenario is to extend the SM by introducing the right-handed Majorana neutrinos with masses around an intermediate scale whose out-of-equilibrium decays create lepton asymmetry in the Universe. The lepton asymmetry is converted to the baryon asymmetry through the $(B + L)$-violating sphaleron transitions [6,7] with the conversion rate [8]

$$Y_B = -\frac{8N_f + 4N_H}{22N_f + 13N_H} Y_L = -\frac{28}{79} Y_L, \tag{2}$$

where we have taken $N_f = 3$ and $N_H = 1$ are the numbers of fermion families and Higgs doublets in the SM. In normal thermal leptogenesis, there is a lower bound on the mass of Majorana neutrinos $\gtrsim 10^{10}$ GeV [9] in order to create sufficient amount of the baryon asymmetry. If it is the case, it is hopeless to directly observe the heavy neutrinos at high-energy colliders in the near future.

Many models beyond the SM have been proposed, which may be realized at the TeV scale and hence accessible to the LHC currently in operation and more future colliders such as the International Linear Collider. Among many models, in this paper, we consider the minimal gauged $B - L$ extended SM. This is an elegant and simple extension of the SM, in which the right-handed neutrinos of three generations are necessarily introduced for the cancellation of the gauge and gravitational anomalies. In addition, the mass of right-handed neutrinos arises associated with the $U(1)_{B-L}$ gauge symmetry breaking and the seesaw mechanism is automatically implemented. In the view point of LHC physics, it is very interesting if the $B - L$ symmetry breaking scale lies around TeV so that the $B - L$ gauge boson (Z' boson) and the right-handed neutrinos can be discovered in the near future [10]. Recently, we have proposed the minimal $B - L$ model with the classical conformal invariance [11] and showed that the $B - L$ symmetry breaking in this model is naturally realized at the TeV scale when the $B - L$ gauge coupling constant is the same order of magnitude as the size of the SM gauge coupling constants [12].

Although the minimal $B - L$ model at TeV is a very attractive scenario, the normal thermal leptogenesis scenario cannot work because the mass scale of the

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right-handed neutrinos is far below the bound, 10^{10} GeV mentioned above. In this case, the CP-asymmetry parameter, which is roughly proportional to Dirac Yukawa coupling squareds is too small to give sufficient amount of baryon asymmetry in the Universe. However, it has been found that when two right-handed neutrinos have almost degenerate masses, there is an enhancement of the CP-asymmetry parameter [13], and this enhancement can make the leptogenesis scenario viable even if the mass scale of the right-handed neutrinos lie around TeV, the resonant leptogenesis [14]. The maximum enhancement is achieved when the mass splitting between two right-handed neutrinos is comparable to the decay width of either right-handed neutrinos. By tuning the mass splitting between two right-handed neutrinos, even a CP-asymmetry parameter of order unity can be obtained in principle. However, it is still nontrivial whether the minimal $B-L$ model at the TeV scale can reproduce the observed baryon asymmetry because, as we will discuss later in detail, the creation of the lepton asymmetry via decays of right-handed neutrinos is highly suppressed in the presence of the $U(1)_{B-L}$ gauge interaction with Z' boson mass at the TeV scale [15].

In this paper, we investigate in detail the resonant leptogenesis scenario in the minimal $B-L$ extended SM with the $B-L$ symmetry breaking at the TeV scale. Through detailed analysis of the Boltzmann equations with a variety of model-parameter sets, we show how much the resultant baryon asymmetry via leptogenesis is enhanced or suppressed, depending on model parameters, in particular, neutrino Dirac Yukawa couplings and TeV-scale Majorana masses of heavy degenerate neutrinos. In order to consider a realistic case, we impose a simple ansatz for model parameters and analyze the neutrino oscillation parameters and the baryon asymmetry via leptogenesis as a function of only a single CP phase. We find that a fixed CP phase can simultaneously reproduce all neutrino oscillation data and the observed baryon asymmetry in the present Universe.

The paper is organized as follows. In the next section, we give a brief review on the minimal $B-L$ model and the natural realization of the $B-L$ symmetry breaking at the TeV scale. In Sec. III, we analyze in detail the resonant leptogenesis at the TeV scale by numerically solving the Boltzmann equations with various parameter sets. We show how the generated baryon asymmetry depends on the model parameters such as Dirac Yukawa coupling, right-handed neutrino mass spectrum, etc. In Sec. IV, we investigate more realistic parameter choices so as to reproduce the neutrino oscillation data. We introduce two right-handed neutrinos and a simple ansatz among the parameters, by which the neutrino oscillation parameters and the baryon asymmetry are determined by only a single CP phase. We find that there exists a CP phase which simultaneously reproduces the neutrino oscillation data and the observed baryon asymmetry. The last section is devoted for conclusions. Formulas used in our analysis are listed in Appendix.

II. THE MINIMAL $B-L$ MODEL AT TEV

The minimal $B-L$ extended SM is based on the gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ with the particle contents listed in Table I. The right-handed neutrinos (N_i) of three generations are necessarily introduced by which all the gauge and gravitational anomalies are canceled. The SM singlet scalar field (Φ) works to break the $U(1)_{B-L}$ gauge symmetry by its vacuum expectation value (VEV), $\langle\Phi\rangle = v_{B-L}/\sqrt{2}$. Once the $B-L$ gauge symmetry is broken, the Z' boson acquires mass,

$$m_{Z'} = 2 g_{B-L} v_{B-L}, \tag{3}$$

where g_{B-L} is the $B-L$ gauge coupling. The current experimental bound was found to be $v_{B-L} \gtrsim 3$ TeV [16].

The Lagrangian relevant for the seesaw mechanism is given by

$$\mathcal{L} \supset -y_D^{ij}\bar{\nu}_R^i H \ell_L^j - \tfrac{1}{2} y_N^i \Phi \bar{\nu}_R^{ic} \nu_R^i + \text{h.c.}, \tag{4}$$

where without loss of generality, we work on the basis in which the second term is diagonalized and y_N^i is real and positive. The first term gives the Dirac neutrino mass term after the electroweak symmetry breaking ($m_D = y_D\langle H\rangle$), while the right-handed neutrino Majorana masses are generated through the second term associated with the $B-L$ gauge symmetry breaking:

$$M_i = \frac{y_N^i}{\sqrt{2}} v_{B-L}. \tag{5}$$

The $B-L$ symmetry breaking scale is determined by parameters in the Higgs potential and in general it can be taken to be any scale as long as the experimental bound $v_{B-L} \gtrsim 3$ TeV [16] is satisfied. As discussed in the previous section, we assume the $B-L$ symmetry breaking at the TeV scale in this paper, and the masses of Z' boson and right-handed neutrinos lie around TeV. In fact, it has been pointed out in [11,12] if we impose the classical conformal

TABLE I. Particle content: In addition to the SM particles, right-handed neutrinos N_i ($i = 1, 2, 3$ denotes the generation index) and a complex scalar Φ are introduced.

	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$
q_L^i	**3**	**2**	$+1/6$	$+1/3$
u_R^i	**3**	**1**	$+2/3$	$+1/3$
d_R^i	**3**	**1**	$-1/3$	$+1/3$
ℓ_L^i	**1**	**2**	$-1/2$	-1
N_i	**1**	**1**	0	-1
e_R^i	**1**	**1**	-1	-1
H	**1**	**2**	$+1/2$	0
Φ	**1**	**1**	0	$+2$

symmetry on the minimal $B-L$ model, the $B-L$ symmetry breaking can be naturally realized at the TeV scale. In the rest of this section, we would like to briefly review the classically conformal $B-L$ extended standard model proposed in [11]. However, since the classical conformal invariance is not important for the leptogenesis scenario (except that it naturally leads to the TeV scale), readers can skip to the next section.

We first note that because of its chiral nature, the SM Lagrangian at the classical level possesses the conformal invariance except for the Higgs mass term, closely related to the gauge hierarchy problem. Bardeen has argued [17] that once the classical conformal invariance and its minimal violation by quantum anomalies are imposed on the SM, it could be free from the quadratic divergences and thus the gauge hierarchy problem. If the mechanism really works, we can directly interpolate the electroweak scale and the Planck scale. Since the classical conformal symmetry forbids the mass term in the Higgs potential, the electroweak symmetry should be broken radiatively through the Coleman-Weinberg (CW) mechanism [18]. Although this is an attractive scenario, the effective Higgs potential is found to be unbounded from below because of the large top Yukawa coupling and therefore the classically conformal SM cannot be a realistic scenario.

In [11], we proposed a classically conformal minimal $B-L$ model and showed that the $B-L$ gauge symmetry breaking is successfully achieved via the CW mechanism and then, this breaking triggers the electroweak symmetry breaking. Because of the CW mechanism, the SM singlet Higgs boson associated with the $B-L$ symmetry breaking is much lighter than Z' boson,

$$\left(\frac{m_\phi}{m_{Z'}}\right)^2 \simeq \frac{6}{\pi}\left(\alpha_{B-L} - \frac{1}{96}\frac{\sum_i(\alpha_N^i)^2}{\alpha_{B-L}}\right) \ll 1, \qquad (6)$$

where $\alpha_{B-L} = g_{B-L}^2/(4\pi)$, and $\alpha_N^i = (y_N^i)^2/(4\pi)$. This formula also indicates the upper bound on α_N^i to keep the vacuum stability, $m_\phi^2 > 0$. Assuming a hierarchical Majorana mass spectrum, for example, we find the upper bound on the heaviest right-handed neutrino mass as

$$\sum_i m_{N_i}^4 < \frac{3}{2} m_{Z'}^4. \qquad (7)$$

Once the $B-L$ symmetry is broken, the Z' boson and the right-handed neutrinos acquire masses at the $B-L$ symmetry breaking scale. Their masses contribute to the effective mass of the SM Higgs doublet through quantum corrections. The naturalness argument, namely, the quantum corrections should not exceed the electroweak scale so far, leads to an upper bound on the $B-L$ symmetry breaking scale. Two-loop corrections with Z' boson and top quarks are found to be dominant, and we have concluded [12] that the $B-L$ symmetry breaking scale should be around TeV.

III. LEPTOGENESIS IN THE MINIMAL $B-L$ MODEL AT TEV

Now we study baryogenesis via leptogenesis in the minimal $B-L$ model at TeV. The lepton asymmetry in the Universe is generated by the CP-violating out-of-equilibrium decays of right-handed neutrinos, and this asymmetry is converted to the baryon asymmetry through the sphaleron process with the conversion rate $Y_B = -(79/28)Y_L$. The generated baryon asymmetry is evaluated by solving the Boltzmann equations. When the Majorana masses of three right-handed neutrinos are largely different as usually assumed, it is sufficient to consider the Boltzmann equations only for the lightest right-handed neutrino. This is because the lepton asymmetry generated by heavier right-handed neutrinos are washed out by the inverse decay of the lightest right-handed neutrinos before its out-of-equilibrium decay [19]. However, in the resonant leptogenesis scenario, (at least) two right-handed neutrinos are degenerate in mass and it is generally not clear whether analysis of the Boltzmann equations with only one right-handed neutrino is sufficient. As we will see later, it can be essential for general cases to consider the Boltzmann equations for multiple right-handed neutrinos.

We begin our analysis with the Boltzmann equations in one-flavor approximation,[1]

$$\begin{aligned}\frac{dY_{N_1}}{dz} &= -\frac{z}{sH(M_1)}\left[\left(\frac{Y_{N_1}}{Y_{N_1}^{\rm eq}}-1\right)(\gamma_{D1}+2\gamma_{h,s}+4\gamma_{h,t}) + \left(\left[\frac{Y_{N_1}}{Y_{N_1}^{\rm eq}}\right]^2-1\right)(\gamma_{Z'}+\gamma_{N,t,\Phi})\right],\\ \frac{dY_{B-L}}{dz} &= -\frac{z}{sH(M_1)}\left[\left(\frac{1}{2}\frac{Y_{B-L}}{Y_l^{\rm eq}}+\epsilon_1\left(\frac{Y_{N_1}}{Y_{N_1}^{\rm eq}}-1\right)\right)\gamma_{D_1} + \frac{Y_{B-L}}{Y_l^{\rm eq}}\left(2(\gamma_N+\gamma_{N,t}+\gamma_{h,t})+\frac{Y_{N_1}}{Y_{N_1}^{eq}}\gamma_{h,s}\right)\right], \end{aligned} \qquad (8)$$

where Y_{N_1} is the yield (the ratio of the number density to the entropy density s) of the (lightest) right-handed neutrino, $Y_{N_1}^{eq}$ is the yield in thermal equilibrium, temperature of the Universe is normalized by the mass of the right-handed neutrino $z = M_1/T$, $H(M_1)$ is the Hubble parameter at $T = M_1$, ϵ_1 is the CP-asymmetry parameter, and γs are the space-time densities of the scatterings in thermal equilibrium. Diagrams in Fig. 1 show processes corresponding to different γs in the Boltzmann equations, whose explicit forms are listed in Appendix. As a good approximation, we neglect masses for all particles involved in the processes, except for the right-handed neutrinos and the Z' boson having the TeV-scale masses. The yield of the right-handed neutrino obeys the first equation, while the send equation determines the $B-L$ number created by the out-of-equilibrium decay of the right-handed

[1] Throughout the paper, our notation follows Ref. [19]

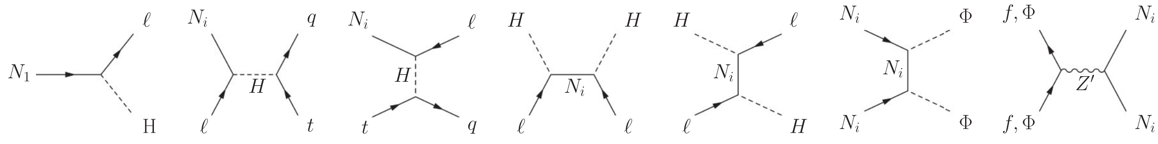

FIG. 1. Feynman diagrams corresponding to each γs.

neutrinos with a nonzero CP-asymmetry parameter. In our numerical studies with input parameters given below, we can check that only γ_{D_1} and $\gamma_{Z'}$ among the space-time densities have important effects on the final results while the others are negligible.

The CP-asymmetry parameter associated with the decay of right-handed neutrino N_i is defined as

$$\epsilon_i \equiv \frac{\sum_j[\Gamma(N_i \to \ell_j H) - \Gamma(N_i \to \ell_j^C H^*)]}{\sum_j[\Gamma(N_i \to \ell_j H) + \Gamma(N_i \to \ell_j^C H^*)]}, \tag{9}$$

which is generated by the interference between the tree and one-loop diagrams shown in Fig. 2. The general formula is given by [13,14]

$$\epsilon_i = -\sum_{j\neq i}\frac{M_i}{M_j}\frac{\Gamma_j}{M_j}\left(\frac{V_j}{2} + S_j\right)\frac{\mathrm{Im}[(y_D y_D^\dagger)_{ij}^2]}{(y_D y_D^\dagger)_{ii}(y_D y_D^\dagger)_{jj}}, \tag{10}$$

where

$$V_j = 2\frac{M_j^2}{M_i^2}\left[\left(1 + \frac{M_j^2}{M_i^2}\right)\log\left(1 + \frac{M_i^2}{M_j^2}\right) - 1\right] \tag{11}$$

corresponds to the vertex correction of the second diagram in Fig. 2 while

$$S_j = \frac{M_j^2 \Delta M_{ij}^2}{(\Delta M_{ij}^2)^2 + M_i^2\Gamma_j^2} \tag{12}$$

is from the self-energy corrections of the third diagram in Fig. 2 with the decay width and the mass difference defined as

$$\frac{\Gamma_j}{M_j} = \frac{(y_D y_D^\dagger)_{jj}}{8\pi} \quad \text{and} \quad \Delta M_{ij}^2 = M_j^2 - M_i^2. \tag{13}$$

In the terminology of $K^0 - \bar{K}^0$ mixing, the first contribution is the so-called direct CP violation while the second one the indirect CP violation. The indirect CP violation occurs because the mass eigenstates and the CP eigenstates are generally different.

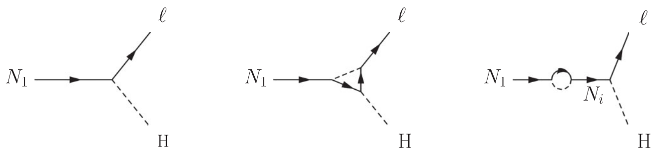

FIG. 2. Right-handed neutrino decay at tree and one-loop levels.

When the right-handed neutrinos have a hierarchical mass spectrum ($M_1 \ll M_{2,3}$), the contributions from V_j and S_j are comparable and the CP-asymmetry parameter is approximately given by [20]

$$\epsilon_1 \sim \frac{3}{16\pi}\frac{m_\nu M_1}{v^2}\sin\delta \sim 10^{-6}\left(\frac{m_\nu}{0.05\ \mathrm{eV}}\right)\left(\frac{M_1}{10^{10}\ \mathrm{GeV}}\right)\sin\delta, \tag{14}$$

where m_ν is the light neutrino mass eigenvalue, and δ is the CP phase and assumed to be of order one. The baryon asymmetry of the Universe is parameterized as

$$Y_B = \kappa\frac{\epsilon_1}{g_*}, \tag{15}$$

where $g_* = \mathcal{O}(100)$ is the number of relativistic degrees of freedom in the early Universe, and κ is the efficiency factor determined by solving the Boltzmann equations, independently of the CP-asymmetry parameter. In the leptogenesis scenario without the Z' gauge boson, this factor is roughly estimated as [9]

$$\kappa \sim 2 \times 10^{-2}\left(\frac{0.05\ \mathrm{eV}}{m_\nu}\right)^{1.1}. \tag{16}$$

Using these estimations, we arrive at the conclusion that the Majorana mass M_1 has to be larger than 10^{10} GeV to give the observed value $Y_B \sim 10^{-10}$. However, note that this conclusion is based on the formula of the CP-asymmetry parameter in the case with the hierarchical right-handed neutrino mass spectrum. In fact, when two right-handed neutrinos are almost degenerate, the CP-asymmetry parameter can be enhanced.

Now, suppose that two right-handed neutrinos are almost degenerate. In this case, it is easy to note that there is a parameter region which can dramatically enhance S_j. Maximum enhancement occurs for $\Delta M_{ij}^2 \simeq M_i\Gamma_j \ll M_i^2$, so that $S_j \sim M_j/\Gamma_j \gg 1$. In this case, the CP-asymmetry parameter is given by

$$\epsilon_i \sim \frac{\mathrm{Im}[(y_D y_D^\dagger)_{ij}^2]}{(y_D y_D^\dagger)_{ii}(y_D y_D^\dagger)_{jj}}, \tag{17}$$

which can, in principle, be of order unity. The leptogenesis scenario with this enhancement of the CP-asymmetry parameter is called the resonant leptogenesis [14]. This enhancement is crucial to realize the observed baryon asymmetry when the right-handed neutrino mass is significantly smaller than 10^{10} GeV, such as the TeV scale which is of our main concern in this paper.

A. Analysis of the Boltzmann equations in one-flavor approximation

Now we analyze the Boltzmann equations in Eq. (8) to see how much baryon asymmetry can be generated in the $B-L$ model at the TeV scale. In order to understand the response between the model parameters involved in this analysis and the resultant baryon asymmetry, we first consider a model in one-flavor approximation with the parameterization of the decay width as

$$\Gamma_1 = \frac{y_D^2}{8\pi} M_1 \tag{18}$$

with a real free parameter y_D, while the other parameters are fixed as follows:

$$\epsilon_1 = 0.01, \qquad \alpha_{B-L} = 0.006, \qquad m_{Z'} = 3\ \text{TeV}, \qquad M_1 = 2\ \text{TeV}. \tag{19}$$

Then, we numerically solve the Boltzmann equations with the boundary conditions

$$Y_{N_1}(0) = Y_{N_1}^{\text{eq}}(0), \qquad Y_{B-L}(0) = 0. \tag{20}$$

The lepton asymmetry generated by the right-handed neutrino decays is converted into the baryon asymmetry via the sphaleron process while the process is in thermal equilibrium. In our analysis throughout the paper, we evaluate the resultant baryon number at the freeze-out temperature of the sphaleron process, $T_{\text{sph}} \simeq 150$ GeV [21], where the conversion of the lepton number to the baryon number is terminated:

$$Y_B = \frac{28}{79} Y_{B-L}(z_{\text{sph}}), \tag{21}$$

where $Y_{B-L}(z_{\text{sph}})$ is the numerical solution of the Boltzmann equations at $z_{\text{sph}} = M_1/T_{\text{sph}}$.

The resultant baryon asymmetry is depicted in Fig. 3 as a function of y_D (solid line). For comparison, results for the cases with $\alpha_{B-L} = 0$ (dotted blue line) and with $\gamma_{N,t,\Phi} = 0$ as well as $\alpha_{B-L} = 0$ (dashed green line) are also shown. For a small $y_D^2 \lesssim 10^{-10.5}$, we can see that the generation of the baryon asymmetry is suppressed in the presence of the Z' boson and $\gamma_{N,t,\Phi}$ processes. Although the suppression by the Z' boson effect dominates, the $\gamma_{N,t,\Phi}$ process mediated by the Majorana Yukawa coupling (y_N) causes a dramatic reduction in generating the baryon asymmetry even in the absence of the Z' boson effect. In the region, Y_B is growing as y_D, and a larger y_D generates a larger baryon asymmetry against the Z' boson and Majorana Yukawa coupling effects. On the other hand, for $y_D^2 \gtrsim 10^{-10}$, the effect by Dirac Yukawa coupling dominates over the Z' boson and Majorana Yukawa coupling effects, and all lines become well-overlapping. In this region, however, Y_B is suppressed by the washing-out process via the inverse-decay process. In the dashed (green) line, Y_B becomes smaller as y_D^2 is lowered for $y_D^2 \lesssim 10^{-15}$ GeV, nevertheless $\gamma_{N,t,\Phi} = 0$ and $\alpha_{B-L} = 0$. This is because the generation of lepton number is too slow with such a small Dirac Yukawa coupling, and the sphaleron process freezes out before the completion of the whole lepton number generation.

Figure 4 shows the results for different values of M_1 as a function of y_D^2 while the other parameters are kept the same. We can see, for $M_1 > 1.5$ TeV, a similar behavior to the result shown in Fig. 3. For a relatively small $M_1 \lesssim 1.5$ TeV, the resultant Y_B becomes almost independent of y_D^2 even for a larger Y_D. This is because the freeze-out of the sphaleron process occurs and thus the conversion of the

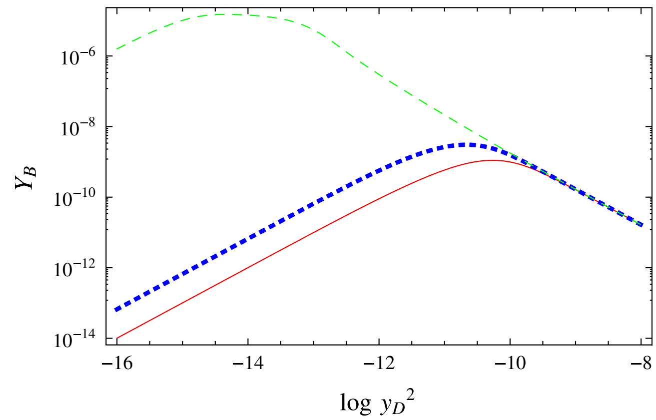

FIG. 3 (color online). The baryon asymmetry generated in the minimal $B-L$ model (solid red line) as a function of the Dirac Yukawa coupling. The dotted (blue) line corresponds to the result of baryon asymmetry in the absence of the $B-L$ gauge interaction, while the dashed (green) line is for the case with $\gamma_{N,t,\Phi} = 0$ as well as $\alpha_{B-L} = 0$.

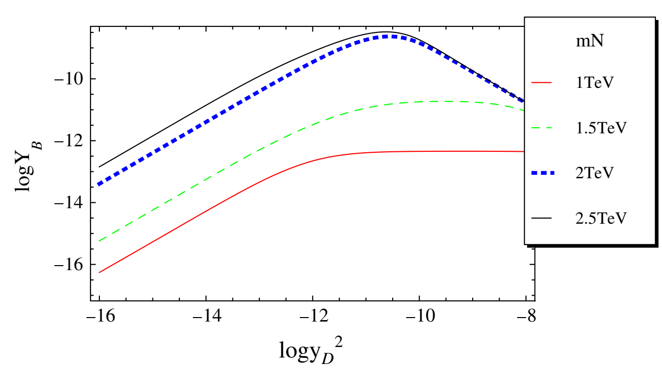

FIG. 4 (color online). The baryon asymmetry as a function of Yukawa coupling for different values of the right-handed neutrino mass, $M_1 = 1$ TeV (solid red line), $M_1 = 1.5$ TeV (dashed green line), $M_1 = 2$ TeV (dotted blue line) and $M_1 = 2.5$ TeV (solid black line).

lepton number to the baryon number is terminated before the washing-out process becomes effective.

B. Two-flavor analysis

We have so far investigated the leptogenesis by solving the Boltzmann equations with one-flavor right-handed neutrino. This treatment is justified when the right-handed neutrinos have a hierarchical mass spectrum. In the resonant leptogenesis, two right-handed neutrinos are degenerated in mass and thus, it is nontrivial whether one-flavor analysis is actually a good approximation or not. Here, we generalize our analysis to the two-flavor case and clarify when the one-flavor analysis is justified.

As we have discussed in the previous section, the CP-asymmetry parameter ϵ_i is maximally enhanced when two right-handed neutrinos are almost degenerate and their mass squared difference is $\Delta M_{ij}^2 \simeq M_i \Gamma_j$. Thus, in our analysis for the two-flavor case, we set the mass difference as $\Delta M_{12}^2 = M_1 \Gamma_2$. Note that since $M_{1,2} \gg \Gamma_{1,2}$ and thus $M_1 \simeq M_2$, the case with $\Delta M_{12}^2 = M_1 \Gamma_1$ is essentially the same as the case with the exchange $1 \leftrightarrow 2$. The CP-asymmetry parameters are given by

$$\epsilon_1 \simeq -\frac{1}{2}\frac{\mathrm{Im}[(y_D y_D^\dagger)_{12}^2]}{(y_D y_D^\dagger)_{11}(y_D y_D^\dagger)_{22}}, \qquad \epsilon_2 \simeq \epsilon_1 \times \frac{2\Gamma_1\Gamma_2}{\Gamma_1^2+\Gamma_2^2}, \tag{22}$$

where we have used the relations $(y_D y_D^\dagger)_{12} = (y_D y_D^\dagger)_{21}^*$ and $\Delta M_{12}^2 = -\Delta M_{21}^2$.

We consider the following three cases:

(1) $\Gamma_1 \gg \Gamma_2$
The CP asymmetry parameter $\epsilon_2 \simeq 2\epsilon_1\Gamma_2/\Gamma_1 \ll \epsilon_1$, and the baryon asymmetry is generated dominantly by the N_1 decay. In addition, the washing-out process by the inverse decay of N_2 is also negligible to that by N_1. Therefore, analysis with only one-flavor right-handed neutrino N_1 is sufficient in evaluating the resultant baryon asymmetry.

(2) $\Gamma_1 \sim \Gamma_2$
Clearly, two right-handed neutrinos are almost identical, so that one-flavor analysis is sufficient, but the resultant baryon asymmetry should be twice of that obtained in the one-flavor case.

(3) $\Gamma_1 \ll \Gamma_2$
The CP asymmetry parameter $\epsilon_2 \simeq 2\epsilon_1 \frac{\Gamma_1}{\Gamma_2} \ll \epsilon_1$ and hence the generation of baryon asymmetry by N_2 decays are negligible. However, the washing-out effect by the inverse decay of N_2 can be efficient and the generated baryon asymmetry can be drastically reduced, depending on the value of Γ_2.

Therefore, for the cases 1 and 2, one-flavor analysis is sufficient to evaluate generated baryon asymmetries, while the case 3 is nontrivial and we need to solve the Boltzmann equations with two-flavor right-handed neutrinos.

The Boltzmann equations with two-flavor right-handed neutrinos are given by

$$\begin{aligned}\frac{dY_{N_1}}{dz} &= -\frac{z}{sH(M_1)}\left[\left(\frac{Y_{N_1}}{Y_{N_1}^{\mathrm{eq}}}-1\right)\gamma_{D_1} + \left(\left(\frac{Y_{N_1}}{Y_{N_1}^{\mathrm{eq}}}\right)^2-1\right)\gamma_{Z'}\right],\\ \frac{dY_{N_2}}{dz} &= -\frac{z}{sH(M_1)}\left[\left(\frac{Y_{N_2}}{Y_{N_2}^{\mathrm{eq}}}-1\right)\gamma_{D_2} + \left(\left(\frac{Y_{N_2}}{Y_{N_2}^{\mathrm{eq}}}\right)^2-1\right)\gamma_{Z'}\right],\\ \frac{dY_{B-L}}{dz} &= -\frac{z}{sH(M_1)}\left[\sum_{j=1}^{2}\left(\frac{1}{2}\frac{Y_{B-L}}{Y_l^{\mathrm{eq}}} + \epsilon_j\left(\frac{Y_{N_j}}{Y_{N_j}^{\mathrm{eq}}}-1\right)\right)\gamma_{D_j}\right].\end{aligned} \tag{23}$$

Here we have omitted γs mediated by the Higgs boson and the right-handed neutrinos, because these effects are, in fact, negligible. However, all processes in Fig. 1 are taken into account in our numerical analysis. With the initial conditions, $Y_{N_i}(0) = Y_{N_i}^{\mathrm{eq}}(0)$ $(i = 1, 2)$ and $Y_{B-L}(0) = 0$, we solve these equations and evaluate the baryon number at $T_{\mathrm{sph}} = 150$ GeV. In order to compare results to those in the one-flavor case, we parameterize the decay width as $\Gamma_1 = y_D^2 M_1/(8\pi)$ corresponding to the one-flavor case. For fixed values of Γ_2/Γ_1, we show the ratio of the baryon asymmetry to the one obtained in the one-flavor analysis, $Y_{B-L}^{2\text{-flavor}}/Y_{B-L}^{1\text{-flavor}}$, in Fig. 5. For a small y_D, we can see that the one-flavor analysis is a good approximation. On the other hand, the washing-out process by the inverse decay of N_2 is very effective for a large y_D, and the baryon asymmetry is very much suppressed than the result obtained in the one-flavor analysis. As is expected, the baryon asymmetry is more suppressed as the ratio Γ_2/Γ_1 becomes larger. Therefore, in order to obtain the correct result for baryon asymmetry via the resonant leptogenesis, analysis with two (or more) flavors can be essential in the general

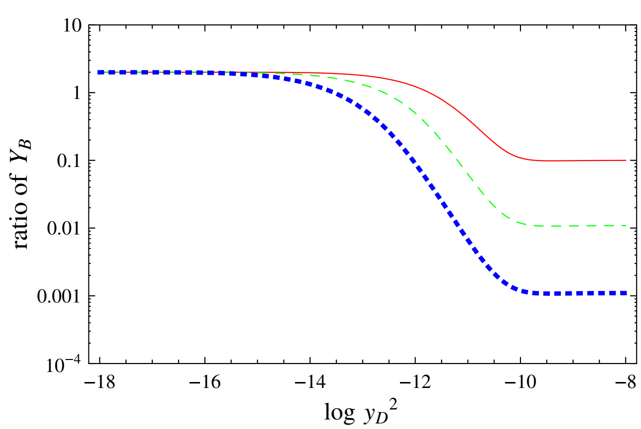

FIG. 5 (color online). The ratio $Y_{B-L}^{2\text{-flavor}}/Y_{B-L}^{1\text{-flavor}}$ as the Dirac Yukawa coupling squared y_D^2, for different choices of $\frac{\Gamma_2}{\Gamma_1} = 10$ (red solid line) $\frac{\Gamma_2}{\Gamma_1} = 100$ (green dashed line) and $\frac{\Gamma_2}{\Gamma_1} = 1000$ (blue dotted line).

case, especially, when $\epsilon_i \ll \epsilon_j$ but $\Gamma_i \gg \Gamma_j$ for two almost degenerate right-handed neutrinos, N_i and N_j.

IV. NEUTRINO OSCILLATION DATA AND RESONANT LEPTOGENESIS

In the previous section, we have analyzed the resonant leptogenesis in the minimal $B-L$ model and investigated the response of the resultant baryon asymmetry to model parameters such as the Dirac Yukawa couplings and the right-handed neutrino masses. It is clearly more interesting to consider a realistic model (in other words, a realistic parameterization) which can account for the observed neutrino oscillation phenomena. For this purpose, we consider a strategy first proposed in [22] in this section. In our analysis, we adopt the current neutrino oscillation data in the 2-σ range [23]:

$$\begin{aligned}
7.25\times 10^{-5} < \Delta m_{12}^2\ (\mathrm{eV}^2) < 8.11\times 10^{-5},\\
2.18\times 10^{-3} < |\Delta m_{13}^2|\ (\mathrm{eV}^2) < 2.64\times 10^{-3},\\
0.27 < \sin^2\theta_{12} < 0.35,\\
0.39 < \sin^2\theta_{23} < 0.63,\\
\sin^2\theta_{13} \le 0.040,
\end{aligned} \tag{24}$$

for the standard parameterization of the mixing matrix,

$$U_{\mathrm{PMNS}} = \begin{pmatrix} \cos\theta_{12}\cos\theta_{13} & \sin\theta_{12}\cos\theta_{13} & \sin\theta_{13}e^{-i\delta} \\ -\sin\theta_{12} - \cos\theta_{12}\sin\theta_{23}\cos\theta_{13}e^{i\delta} & \cos\theta_{12}\cos\theta_{23} - \sin\theta_{12}\sin\theta_{23}\sin\theta_{13}e^{i\delta} & \sin\theta_{23}\cos\theta_{13} \\ \sin\theta_{12}\sin\theta_{23} - \cos\theta_{12}\cos\theta_{23}\sin\theta_{13}e^{i\delta} & -\cos\theta_{12}\cos\theta_{23} - \sin\theta_{12}\cos\theta_{23}\sin\theta_{13}e^{i\delta} & \cos\theta_{23}\cos\theta_{13} \end{pmatrix} \times \mathrm{diag}(e^{i(\alpha_1/2)}, e^{i(\alpha_2/2)}, 1), \tag{25}$$

with the Dirac phase δ and the Majorana phases α_i.

We consider the so-called minimal seesaw [24] in the context of the minimal $B-L$ model, and assume that only two right-handed neutrinos are relevant for the neutrino oscillation phenomena and leptogenesis. The third right-handed neutrino is assumed to decouple from the neutrino oscillation phenomena by some reason. A simple idea is to introduce a discrete Z_2 symmetry under which the third right-handed neutrino is assigned to be odd while all the other particles in the $B-L$ model even. In the context of the minimal $B-L$ model with the Z_2 symmetry, it has been shown [25] that the third right-handed neutrino can be a suitable candidate for the cold dark matter with the relic density consistent with observations.

In the minimal seesaw model, we parameterize the 2×3 Dirac neutrino mass matrix, without loss of generality, as

$$m_D = \begin{pmatrix} a_1 e^{i\phi_1} & a_2 e^{i\phi_2} & a_3 e^{i\phi_3} \\ a_4 & a_5 & a_6 \end{pmatrix}, \tag{26}$$

where a_i and ϕ_j are real parameters, and we have worked in the basis where both the charged lepton mass matrix and the right-handed neutrino mass matrix are diagonalized with real and positive eigenvalues. We parameterize the Majorana mass matrix of the two right-handed neutrinos as

$$M_N = \begin{pmatrix} M_1 & 0 \\ 0 & M_1(1+r) \end{pmatrix}, \tag{27}$$

where the parameter r should be very small, for example, $r \sim \Gamma_1/M_1$ or Γ_2/M_1 in order to realize the enhancement of the CP-asymmetry parameter. Although r is crucial for the resonant leptogenesis, such a small r is negligible in fitting for the neutrino oscillation data.

For simplicity, we fix $\phi_1 = \phi_2 = 0$ in our analysis and introduce an ansatz [22] that the light neutrino mass matrix after the seesaw mechanism [5],

$$m_\nu = m_D^T M_N^{-1} m_D \simeq \frac{1}{M_1} m_D^T m_D, \tag{28}$$

is diagonalized by the so-called tribimaximal mixing matrix [26],

$$U_{TB} = \begin{pmatrix} \sqrt{\frac{2}{3}} & \sqrt{\frac{1}{3}} & 0 \\ -\sqrt{\frac{1}{6}} & \sqrt{\frac{1}{3}} & \sqrt{\frac{1}{2}} \\ -\sqrt{\frac{1}{6}} & \sqrt{\frac{1}{3}} & -\sqrt{\frac{1}{2}} \end{pmatrix}, \tag{29}$$

in the CP invariant case ($\phi_3 = 0$). As is well-known, this tribimaximal mixing matrix gives almost the best fit in the oscillation data.

Note that in the minimal seesaw, the rank of the light neutrino mass matrix is two and the lightest mass eigenvalue is 0. In the following we consider two cases for the light neutrino mass spectrum, namely, the normal hierarchical (NH) case and inverted-hierarchical (IH) case. Let us first consider the NH case, where we have

$$D_\nu^{\mathrm{NH}} = \mathrm{diag}(0, m_2^{\mathrm{NH}}, m_3^{\mathrm{NH}}) \tag{30}$$

with $m_2^{\mathrm{NH}} = \sqrt{\Delta m_{12}^2}$ and $m_3^{\mathrm{NH}} = \sqrt{|\Delta m_{13}^2|}$. According to our ansatz, we first find a solution to $m_\nu = U_{\mathrm{TB}} D_\nu^{\mathrm{NH}} U_{\mathrm{TB}}^T$ in the CP-invariant case. Among several solutions, we choose, as an example,

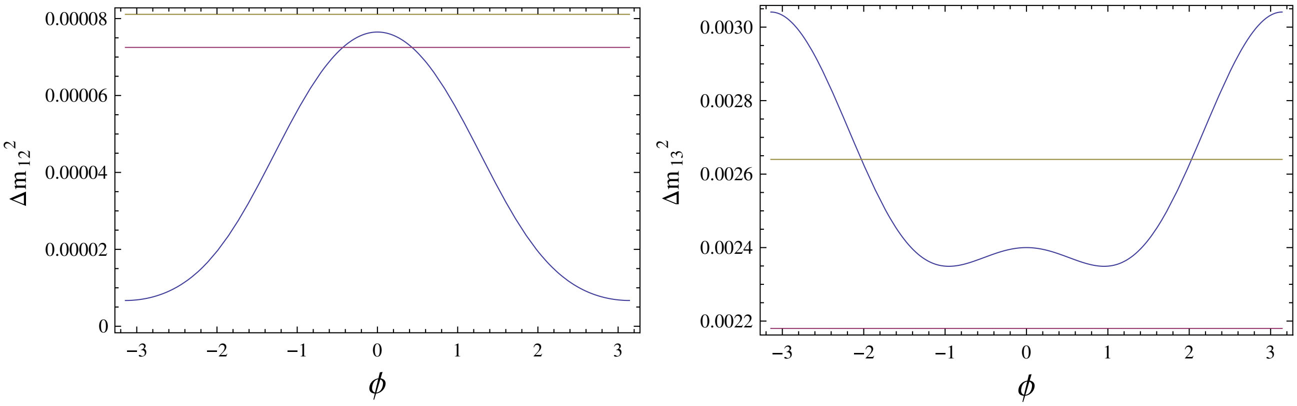

FIG. 6 (color online). The neutrino oscillation parameters, Δm_{12} (left panel) and Δm_{13} (right panel), as a function of the CP phase ϕ_3. The observed data in the 2-σ range are indicated by two horizontal lines.

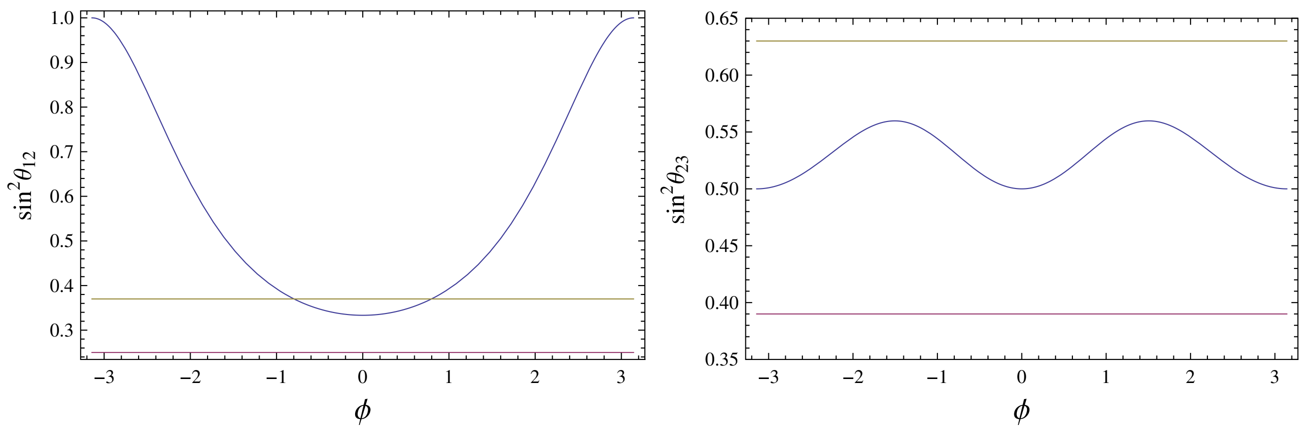

FIG. 7 (color online). The neutrino oscillation parameters, $\sin^2\theta_{12}$ (left panel) and $\sin^2\theta_{23}$ (right panel), as a function of the CP phase ϕ_3. The observed data in the 2-σ range are indicated by two horizontal lines.

$$a_1 = a_2 = a_3 = \sqrt{\frac{M_1 m_2^{\rm NH}}{3}},$$
$$a_4 = 0, \qquad (31)$$
$$a_5 = -a_6 = \sqrt{\frac{M_1 m_3^{\rm NH}}{2}}.$$

For the input values, we use $m_2^{\rm NH} = 8.75 \times 10^{-3}$ eV and $m_3^{\rm NH} = 4.90 \times 10^{-2}$ eV. In this section, we fix other parameters as $\alpha_{B-L} = 0.006$, $m_{Z'} = 3$ TeV and $M_1 = 2$ TeV.

Now we turn the CP phase ϕ_3 on. With fixed a_i and M_1, the light neutrino mass matrix is given as a function of the single parameter ϕ_3 [22] (in the approximation with $r = 0$). In Figs. 6–8, the neutrino oscillation parameters are depicted as a function of the CP phase ϕ_3. As we expect, the outputs of the oscillation parameters deviate from the values at $\phi_3 = 0$ as the CP phase is changed, and eventually some of outputs are found to be outside of the 2-σ range of the experimental data. We find the bound on the CP phase as $|\phi_3| \lesssim 0.5$.

For the resonant leptogenesis, both $\phi_3 \neq 0$ and $r \neq 0$ are crucial. Figure 9 shows the CP-asymmetry parameters (ϵ_1 and ϵ_2) as a function of r with $\phi_3 = 0.5$ for example. For the right (left) curve corresponding to ϵ_1 (ϵ_2), a peak appears around $r = 10^{-13}$ ($r = 10^{-14}$). Choosing

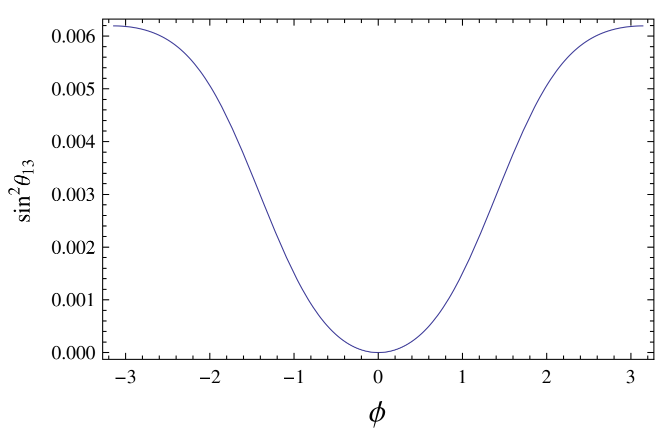

FIG. 8 (color online). $\sin^2\theta_{13}$ as a function of the CP phase.

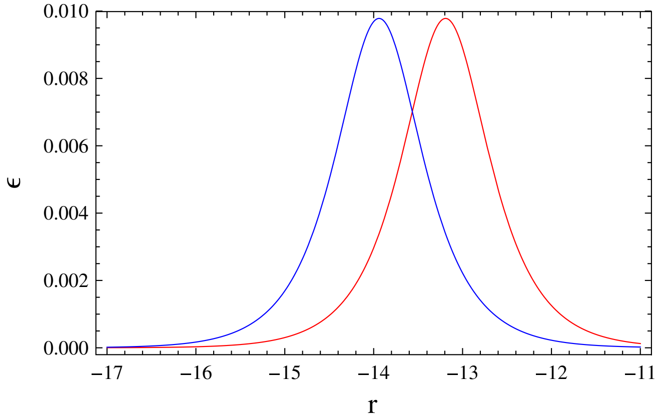

FIG. 9 (color online). The CP asymmetry parameters, ϵ_1 (right) and ϵ_2 (left), as a function of r.

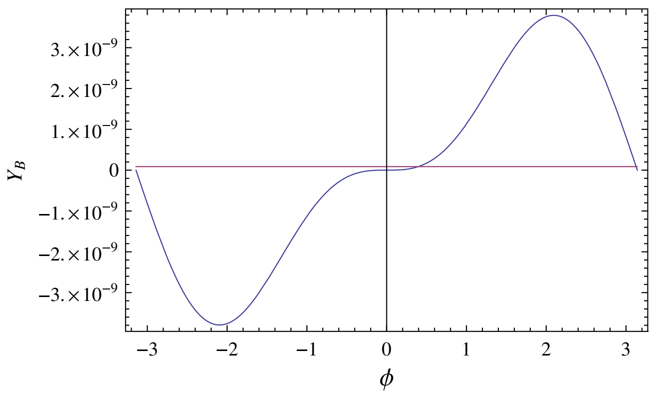

FIG. 10 (color online). The resultant baryon asymmetry in the Universe as a function of the CP phase ϕ_3. The observed value $Y_B = 0.87 \times 10^{-10}$ is depicted as the horizontal line.

$r = 10^{-14}$ for example, the CP-asymmetry parameters are also given as a function of only ϕ_3. Therefore, we have correlations between neutrino oscillation parameters and the baryon asymmetry generated by the resonant leptogenesis through the CP phase ϕ_3 [22]. Interestingly, the amount of the generated baryon asymmetry becomes larger as ϕ_3 goes away from zero, while a large displacement of ϕ_3 from zero results in the oscillation parameters inconsistent with the experimental data.

Numerical solution of the Boltzmann equations with two flavors in Eq. (23) is shown Fig. 10 as a function of ϕ_3. We find that the observed baryon asymmetry $Y_B = 0.87 \times 10^{-10}$ in the present Universe is obtained for $\phi_3 = 0.35$, for which the neutrino oscillation parameters are fixed as

$$\begin{aligned} \Delta m_{12}^2\ (\mathrm{eV}^2) &= 7.39 \times 10^{-5}, \\ \Delta m_{13}^2\ (\mathrm{eV}^2) &= 2.39 \times 10^{-3}, \\ \sin^2\theta_{12} &= 0.34, \\ \sin^2\theta_{23} &= 0.51, \\ \sin^2\theta_{13} &= 0.00016. \end{aligned} \tag{32}$$

They are all consistent with observations. Although a nonvanishing $\sin^2\theta_{13}$ is predicted, it is quite small, far below the current upper bound.

Next we consider the IH case, where the light neutrino mass matrix is diagonalized as

$$D_\nu^{\mathrm{IH}} = \mathrm{diag}(m_1^{\mathrm{IH}}, m_2^{\mathrm{IH}}, 0), \tag{33}$$

with $m_1^{\mathrm{IH}} = \sqrt{|\Delta m_{13}^2|}$ and $m_2^{\mathrm{IH}} = \sqrt{\Delta m_{12}^2 + |\Delta m_{13}^2|}$. In the CP invariant case, we choose a solution to $m_\nu = U_{\mathrm{TB}} D_\nu^{\mathrm{IH}} U_{\mathrm{TB}}^T$ as

$$\begin{aligned} a_1 = a_2 = a_3 &= \sqrt{\frac{M_1 m_2^{\mathrm{IH}}}{3}}, \\ a_4 &= \sqrt{\frac{2M_1 m_1^{\mathrm{IH}}}{3}}, \\ a_5 = a_6 &= -\sqrt{\frac{M_0 m_{1,\mathrm{IH}}}{6}}. \end{aligned} \tag{34}$$

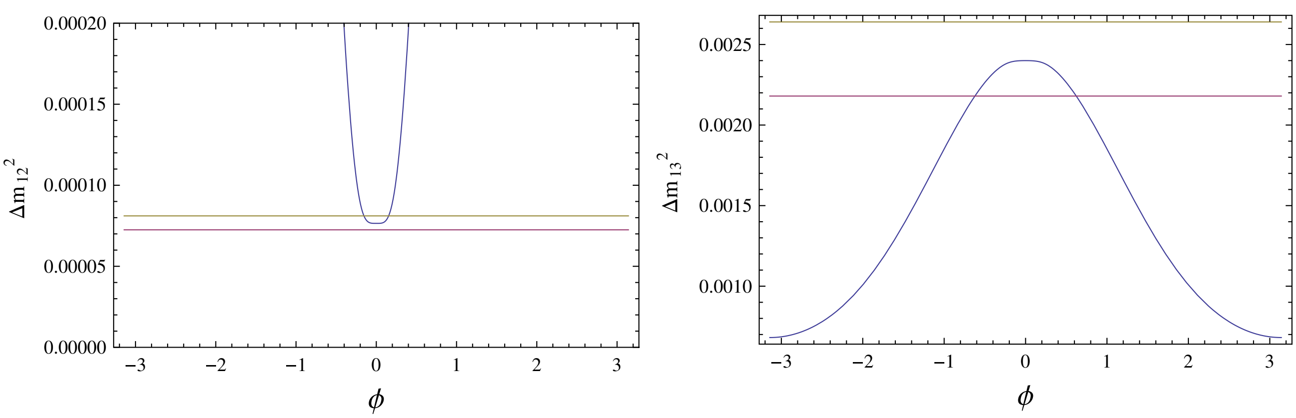

FIG. 11 (color online). The neutrino oscillation parameters, Δm_{12} (left panel) and Δm_{13} (right panel), as a function of the CP phase ϕ_3. The observed data in the 2-σ range are indicated by two horizontal lines.

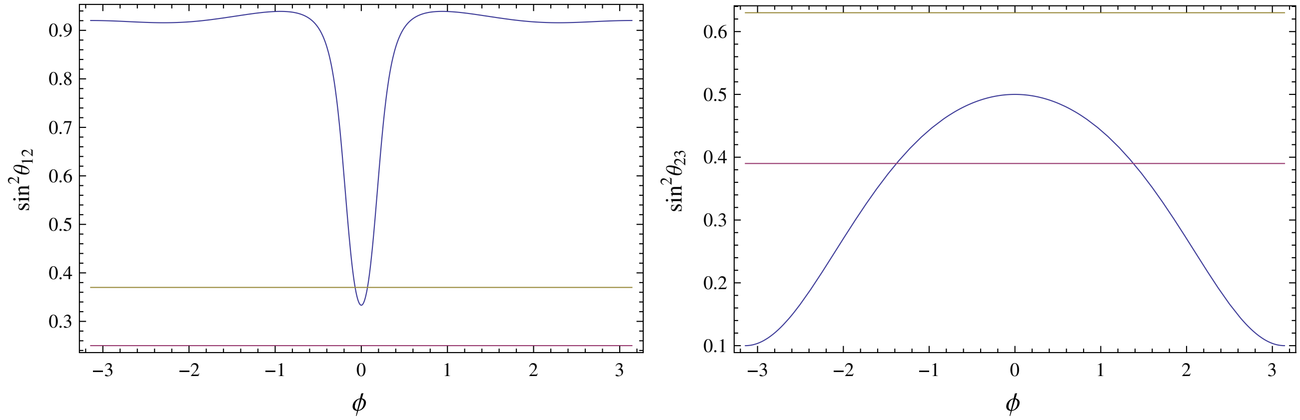

FIG. 12 (color online). The neutrino oscillation parameters, $\sin^2\theta_{12}$ (left panel) and $\sin^2\theta_{23}$ (right panel), as a function of the *CP* phase ϕ_3. The observed data in the 2-σ range are indicated by two horizontal lines.

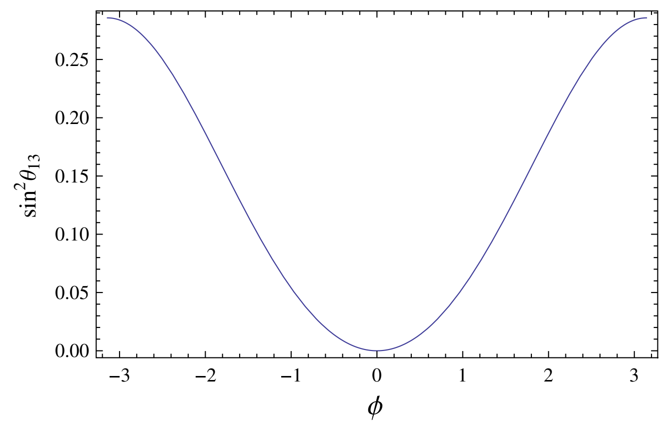

FIG. 13 (color online). $\sin^2\theta_{13}$ as a function of the *CP* phase ϕ_3.

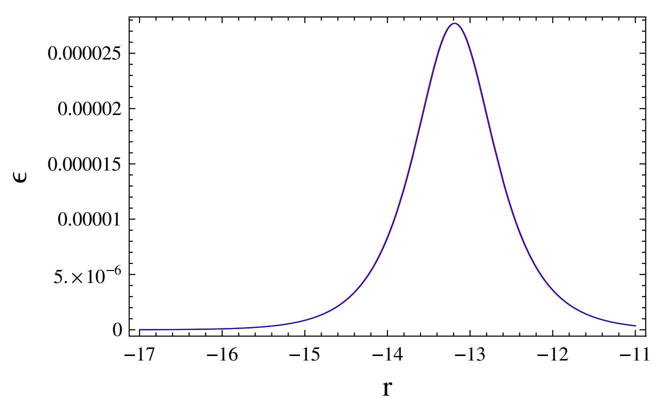

FIG. 14 (color online). The *CP* asymmetry parameters, ϵ_1 and ϵ_2 as a function of r. The two curves are well overlapped.

For the input values of $m_1^{\rm IH} = 4.90 \times 10^{-2}$ eV and $m_2^{\rm IH} = 4.98 \times 10^{-2}$ eV, the neutrino oscillation parameters are depicted in Figs. 11–13 as a function of ϕ_3. A *CP* phase $|\phi_3| \lesssim 0.1$ results in the outputs of the neutrino oscillation

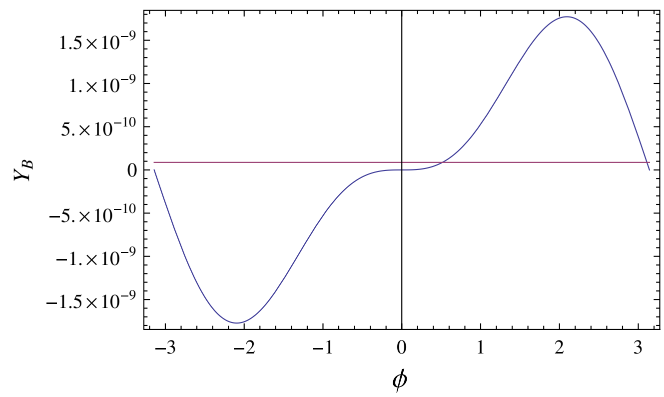

FIG. 15 (color online). The resultant baryon asymmetry as a function of the *CP* phase ϕ_3, along with the observed value $Y_B = 0.87 \times 10^{-10}$ (horizontal line).

parameters consistent with the experimental data in the 2-σ range.

Figure 14 shows the *CP*-asymmetry parameters, ϵ_1 and ϵ_2, as a function of r for $\phi_3 = 0.1$. Two curves are well-overlapped and the peak appears around $r \simeq 10^{-13}$. Then, we solve the Boltzmann equations for $r = 10^{-13}$, and show the results in Fig. 15. Although the observed baryon asymmetry in the Universe is generated for $\phi_3 = 0.43 > 0.1$, the neutrino oscillation parameters corresponding to the *CP* phase are outside of the 2-σ range. Therefore, in the present scheme, the IH case cannot reproduce the neutrino oscillation data and the observed baryon asymmetry simultaneously.

V. CONCLUSIONS

In this paper, we have investigated a possibility to explain the baryon asymmetry of the Universe as well as the neutrino oscillation data in the TeV scale $B - L$ model. In the model, the lepton asymmetry is generated in the early Universe via out-of-equilibrium decays of the right-handed neutrinos with the CP-asymmetry parameter, and converted into the baryon asymmetry via the sphaleron process. When the mass scale of the right-handed neutrinos is low $\lesssim 10^{10}$ GeV, the enhancement of the CP-asymmetry parameter is crucial in order to generate sufficient amount of baryon asymmetry in the Universe. The enhancement is realized when two right-handed neutrinos are almost degenerated and in this case, the CP-asymmetry parameter can be in principle order unity. This scenario is called the resonant leptogenesis. However, it is still nontrivial whether the resonant leptogenesis can realize the observed baryon asymmetry in the context of the minimal $B-L$ model, because the $B-L$ interaction mediated by the Z' boson can dramatically reduce the generation of baryon asymmetry.

We numerically solved the Boltzmann equations for the resonant leptogenesis in the minimal $B-L$ model, and figured out the response between the generated baryon asymmetry and the model parameters such as the neutrino Dirac Yukawa couplings and the right-handed neutrino masses. We first analyzed the Boltzmann equations with only one-flavor right-handed neutrino and a fixed CP-asymmetry parameter. When the neutrino Dirac Yukawa coupling is small, $y_D \lesssim 10^{-10.5}$, the amount of the baryon asymmetry becomes larger as the Dirac Yukawa coupling is raised. In this parameter region, the Z' boson and Majorana Yukawa coupling effects dramatically suppress the generation of baryon asymmetry. For a large Dirac Yukawa coupling, the baryon number generation by the right-handed neutrino decay dominates over the suppression by the Z' boson effect. However, a too large Dirac Yukawa coupling in turn suppresses the generation of the baryon asymmetry by the washing-out effect via the inverse-decay process. Next, we have analyzed the Boltzmann equations with two-flavor right-handed neutrinos and shown that two-flavor analysis can be essential in general cases. With these analyses, we have shown that in some areas of the parameter space a sufficient amount of the baryon asymmetry can be generated though the resonant leptogenesis in the TeV-scale $B-L$ model.

Finally, we have checked whether these parameters are consistent with the current neutrino oscillation data. We have introduced a simple ansatz for the neutrino mass matrices, by which the neutrino Dirac Yukawa couplings are determined as a function of a single CP phase. For both the normal hierarchical and inverted-hierarchical mass spectra of the light neutrinos, we have shown the correlations between the neutrino oscillation parameters and the generated baryon asymmetry via the resonant leptogenesis. In our analysis with the ansatz, a fixed CP phase can reproduce simultaneously the neutrino oscillation data and the observed baryon asymmetry in the normal hierarchical case. On the other hand, we cannot find such a CP phase in the inverted-hierarchical case.

ACKNOWLEDGMENTS

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APPENDIX

The number density n_ψ of a particle ψ (a right-handed neutrino in our case) with mass m_ψ in the early Universe is evaluated by solving the Boltzmann equation of the form [27],

$$\frac{dY_\psi}{dz} = -\frac{z}{sH(m_\psi)} \sum_{a,i,j,\ldots} \bigg[\frac{Y_\psi Y_a \cdots}{Y_\psi^{\rm eq} Y_a^{\rm eq} \cdots} \times \gamma^{\rm eq}(\psi + a + \cdots \to i + j + \cdots) - \frac{Y_i Y_j \cdots}{Y_i^{\rm eq} Y_j^{\rm eq} \cdots} \gamma^{\rm eq}(i + j + \cdots \to \psi + a + \cdots) \bigg], \tag{A1}$$

where $Y_\psi = n_\psi / s$ is the ratio of n_ψ and the entropy density s, $z = \frac{m_\psi}{T}$, and $H(m_\psi)$ is the Hubble parameter at a temperature $T = m_\psi$. The right hand side of Eq. (A1) describes the interactions that change number of ψ, and $\gamma^{\rm eq}$ is the space-time density of scatterings in thermal equilibrium. For a dilute gas we take into account decays, two-particle scatterings and the corresponding back reactions. One finds, for a decay the particle ψ,

$$\gamma_D = \gamma^{\rm eq}(\psi \to i + j + \cdots) = n_\psi^{\rm eq} \frac{K_1(z)}{K_2(z)} \tilde{\Gamma}_{\rm rs}, \tag{A2}$$

where K_1 and K_2 are the modified Bessel functions, and $\tilde{\Gamma}_{\rm rs}$ is the decay width. For two body scattering one has

$$\gamma(\psi + a \leftrightarrow i + j + \cdots) = \frac{T}{64\pi^4} \int_{(m_\psi + m_a)^2}^{\infty} \times ds \hat{\sigma}(s) \sqrt{s} K_1\left(\frac{\sqrt{s}}{T}\right), \tag{A3}$$

where s is the squared center-of-mass energy and the reduced cross section $\hat{\sigma}$ for the process $\psi + a \leftrightarrow i + j + \cdots$ is related to the usual total cross section $\sigma(s)$ by

$$\hat{\sigma}(s) = \frac{8}{s}[(p_\psi \cdot p_a)^2 - m_\psi^2 m_a^2]\sigma(s). \tag{A4}$$

In the following we list the explicit forms of the reduced cross sections used in our analysis [19]. The reduced cross section corresponding to γ_N is given by

$$\hat{\sigma}_N(s) = \frac{\alpha^2}{\sin^4\theta}\frac{2\pi}{M_W^4}\frac{1}{x}\left[a_1(m_D m_D^\dagger)_{11}^2\left(x + \frac{2x}{D_1(x)} + \frac{x^2}{2D_1^2(x)} - \left(1 + 2\frac{x+1}{D_1(x)}\right)\ln(x+1)\right)\right], \tag{A5}$$

where $x = \frac{s}{m_N^2}$, and

$$D_1(x) = x - 1 + \frac{c}{x-1}, \quad \text{with} \quad c = \left(\frac{\tilde{\Gamma}_{\rm rs}}{m_N}\right)^2, \tag{A6}$$

while the one corresponding to $\gamma_{N,t}$ is

$$\hat{\sigma}_{N,t}(s) = \frac{2\pi\alpha^2}{M_W^4\sin^4\theta}\left[(m_D m_D^\dagger)_{11}^2\left(\frac{x}{2(x+1)} + \frac{1}{x+2}\ln(x+1)\right)\right]. \tag{A7}$$

The reduced cross section for the t-channel (and u-channel) process $N + N \to \Phi + \Phi$ mediated by the right-handed neutrino is given by

$$\hat{\sigma}_{N,t,\Phi}(s) = \frac{y_N^4}{8\pi}\frac{x-4}{x}\left(-2 + \frac{x}{2} + \frac{x^2 - 8x + 16}{x\sqrt{x(x-4)}} \times \log\frac{x - \sqrt{x(x-4)}}{x + \sqrt{x(x-4)}}\right). \tag{A8}$$

The reduced cross section for the s-channel process $N + l \to \bar{t} + q$ mediated by the Higgs doublet is given by

$$\hat{\sigma}_{h,s}(s) = \frac{3\pi\alpha^2 m_t^2}{M_W^4\sin^4\theta}(m_D m_D^\dagger)_{11}\left(\frac{x-1}{x}\right)^2, \tag{A9}$$

while for the t-channel process

$$\hat{\sigma}_{h,t}(s) = \frac{3\pi\alpha^2 m_t^2}{M_W^4\sin^4\theta}(m_D m_D^\dagger)_{11} \times \left[\frac{x-1}{x} + \frac{1}{x}\ln\left(\frac{x-1+y'}{y'}\right)\right] \tag{A10}$$

with $y' = \frac{m_h^2}{M_1^2}$. The total reduced cross section for the process $f + \bar{f}, \Phi + \Phi \to N + N$ mediated by the Z' boson (f denotes the SM fermions) is given by

$$\hat{\sigma}_{Z'}(s) = \frac{104\pi}{3}\alpha_{B-L}^2\frac{\sqrt{x}}{(x-y)^2 + yc}(x-4)^{3/2}, \tag{A11}$$

where $y = \frac{m_{Z'}^2}{m_N^2}$, and $c = (\tilde{\Gamma}_{Z'}/M_1)^2$ with the decay width of Z' boson

$$\tilde{\Gamma}_{Z'} = \frac{\alpha_{B-L}m_{Z'}}{6}\left[3\left(1 - \frac{4}{y}\right)^{3/2}\theta(y-4) + 13\right]. \tag{A12}$$

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